Impact of Coulomb interaction and Kondo effect on transport in quantum dots

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We examine the impact of Coulomb electron-electron interaction on transport in a junction with a quantum dot described by Kondo Hamiltonian. We analyze the Fermi liquid regime and consider the limit of zero temperature. With the help of Keldysh technique we calculate the non-linear current and shot noise as a function of applied voltage. We show that Coulomb interaction markable influences the universal effective charge of current-carrying particles $e^* = 5/3e$ which can be measured in shot-noise experiments. The electron-electron interaction modifies this universal value by a factor ($e^* = 5/3eF$) which is less then one and also voltage dependent.

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The non-linear conductance and zero-frequency noise out of equilibrium (shot noise) can yield information on charge fluctuations in physical systems [1]. Indeed, in many cases the effective charge of carriers may be different from bare electron charge e. This occurs when the fundamental quasiparticle charge does not coincide with e or due to interactions. The Kondo effect in the quantum dots opens the opportunity to investigate the changing of effective charge due to interactions. For this purpose it is of a special interest the Fermi liquid regime of Kondo model. In this case the backscattering current is defined by the scattering term and by the interaction quadratic in the spin current. The measurements of the backscattering current and shot noise can give the information about effective backscattering charge of the carriers influenced by interactions. These reasons have been put forward in a recent work [1] where the effective current-carrying charge of the particles was found to acquire the universal value $e^*/e = 5/3$. However, as we will show in the present article, the Coulomb interaction and corresponding electrical potential fluctuations have a noticeable impact on this effective charge. They change the value of backscattering charge and cause a voltage dependence of e^*. To obtain the effective charge with or without electrical potential fluctuations we have to find both parts (scattering and interaction) of the current and shot noise [1, 2, 3]. We concentrate in the following on $T = 0$ limit, where the thermal noise vanishes. We consider the regime around the strongly interacting fix point (Fermi liquid regime) of Kondo Hamiltonian, that is, $eV << T_K$ (T_K is the Kondo temperature) counting for potential fluctuation at the one loop approximation. Here we notice that the role of Coulomb interaction in the case of a large dots was analyzed before in a number of articles [4, 5, 6].

Our principal result can be written compactly as $e^*/e = \frac{5}{3}F$, where factor F depends on the bias voltage, charging energy $E_C = e^2/C$ (here C is the total capacitance), and the effective number of conducting channels $\bar{N}$. In Fig.1 we plot this factor as a function of voltage

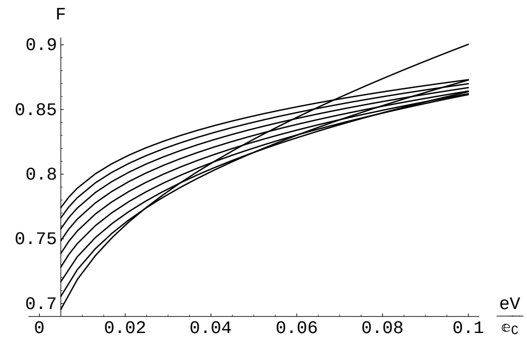

FIG. 1: Voltage dependence (in units of Coulomb energy E_C) of factor F. The curves are related to different number of effective conducting channels $\bar{N}$ starting from $\bar{N}$=2 up to $\bar{N}$=10

for different $\bar{N}$. For validity of perturbation theory that we use, the absolute value of the voltage must be not too small. Also for dot systems $E_C >> T_K$. In this limit we have

$$F = \frac{1 + a_n}{1 + a_j} \quad (1)$$

$$a_n = \frac{4}{5\bar{N}}(\ln(\frac{2E_C\bar{N}}{\pi eV} + 0.36) \quad (2)$$

$$a_j = \frac{7}{4\bar{N}}(\ln(\frac{2E_C\bar{N}}{\pi eV} + 0.16) \quad (3)$$

In the region of small voltages factor F (Fig1.) rather strongly (30%) reduces the effective charge. Now we present the basic steps which lead to the derivation of equation (1).

Effective action: Coulomb interaction in a quantum dot is responsible for formation of Kondo effect and for electric potential fluctuations in the whole junction. The Kondo effect depends also on the tunnelling rates between the dot and the leads. The later strongly influence the Kondo temperature T_K. The Coulomb interaction is characterized by the charging energy E_C. The parameter that controls the impact of electron electron interactions

on the tunnelling is the total conductance g or the number of open channels N. More precisely, in the case of large N we can find the interaction corrections by performing $1/N$ expansion. Here we address the unitary limit of Kondo quantum dot with odd numbers of electrons in it, that is, we deal with ordinary Kondo effect in a quantum dot. Though, in this case $N = 2$ (effective N may be increased by a shunting resistance) we believe that even at $1/N$ level the qualitative picture is correct. The problem is also interesting because of its relation to a rare case of a non-equilibrium system with interaction where the electrical potential fluctuations are included. As an advantage of the unitary limit is that the effective transmissions are close to the unity. Indeed, conductance in additions to $\sigma_0 = 2e^2/h$ has a small nonlinear contribution $\sigma_V \sim -V^2/T_K^2$, where the bias voltage $V << T_K$. As we will show the Coulomb interaction modifies only the nonlinear part of the transport current and shot noise. Therefore, to the lowest order in V/T_K one can calculate the effective action for electric potentials taking in considerations only the unitary limit. In this limit we can use two approaches which give the same result for spin 1/2 Kondo model: one is based on scattering representation of the Kondo Hamiltonian [8], the other is the mean field slave boson approximation (SB) [9]. The SB is more transparent and it will be used to obtain the effective action. We start with the Anderson Hamiltonian

$$\begin{aligned}\hat{H} &= H_L + H_R + \sum_{k,\sigma,\alpha}(v_\alpha c^\dagger_{\alpha\sigma,k} d_\sigma + \text{H.c.}) \\ &+ \sum_\sigma \epsilon d^\dagger_\sigma d_\sigma + U d^\dagger_\uparrow d_\uparrow d^\dagger_\downarrow d_\downarrow \end{aligned} \tag{4}$$

The first two terms correspond to non-interacting electrons in the two leads

$$H_{L(R)} = \sum_{k,\sigma} \xi_{L(R)k} c^\dagger_{L(R)\sigma,k} c_{L(R)\sigma,k} \tag{5}$$

where $c_{\alpha\sigma,k}$, $\xi_{\alpha k}$ are the electron field operator and the electron energy of a lead. Index $\alpha = L, R$ indicates left (right) lead. We assume that the leads are biased by the fluctuating electric field potential drop $v(t) = \frac{d\varphi(t)}{dt}$. There are also electron electron interactions in the leads which are not included in the Hubbard coupling U. At first we separate the Hubbard interaction that leads to Kondo effect. For this we apply to Hamiltonian (4) the SB approximation at $U \to \infty$. In the slave-boson approach, the localized electron operator $d^\dagger_\sigma$ is represented by $\hat{f}^\dagger_\sigma \hat{b}$ with $\hat{b}$ and $\hat{f}^\dagger_\sigma$ being, respectively, the standard boson and fermion operators. The action on Keldysh contour then will incorporate the Anderson Hamiltonian (4) with the new boson and fermion operators. The quadric term is replaced by the two new terms: one represents the renormalized level position in the dot $\tilde{\epsilon}$, and the second stands for requirement of a single occupancy which demands a Lagrange multiplier λ

$$\hat{b}^\dagger \hat{b} + \sum_\sigma \hat{f}^\dagger_\sigma \hat{f}_\sigma = 1 \tag{6}$$

In the SB approximation the Bose operators $\hat{b}^\dagger, \hat{b}$ are replaced by their expectation value b. Having the Anderson part of the action arranged we can get the general form of the total effective action by including the Coulomb interaction term. Performing the standard Hubbard-Stratonovich decoupling of this term in the action, one introduces on the Keldysh contour fluctuating electrical potentials $v_i(t) = \frac{d\varphi_i(t)}{dt}, (i = 1, 2)$ [4, 10]. Next we integrate out the fermion operators in the leads, and also add the action of external circuit expressed in terms of shunt resistance R_s. Thus the effective action acquires a form $S_{eff} = S_b + S + S_s$, where nonoperator bosonic part of the action is

$$S_b = -\int \lambda_i (b_i^2 - 1)\sigma_z^{ii} \tag{7}$$

where σ_z is the Pauli matrix. The action S consists of two parts: the one describing the scattering of electron in the leads, and the other (S_{HS}) responsible for the fluctuating Habbard-Stratonovich fields

$$S = -iTrLnG^{-1} + S_{HS} \tag{8}$$

$$G^{-1}(t,t') = [i\frac{\partial}{\partial t} - \tilde{\epsilon} - \frac{\Gamma b^2}{2\pi}\sigma_z(\tilde{g}_L + \tilde{g}_R)]\sigma_z \tag{9}$$

$$\tilde{g}_{L,R}(t,t') = e^{\pm i\hat{\varphi}(t)} g_{L,R}(t-t') e^{\mp i\hat{\varphi}(t')}, \tag{10}$$

$$S_{HS} = \frac{C}{2e^2}\int dt(tr\dot{\hat{\varphi}}\sigma_z\dot{\hat{\varphi}}) \tag{11}$$

$$\begin{aligned} S_s &= \frac{1}{e^2 R_s}[\frac{i}{2}\int\int dtdt'\chi(t)\gamma(t-t')\chi(t') \\ &+ \int dt\chi(t)(eV - (\dot{\varphi}_1 + \dot{\varphi}_2)/2)] \end{aligned} \tag{12}$$

here Γ is a tunnelling width, $\chi = \varphi_1 - \varphi_2$, the Fourier transform of $\gamma(t)$ at finite temperature T is $\gamma(\omega) = \omega coth(\omega/2T)$. Here and below the values with the hat $\hat{v}, \hat{\varphi}$ mean the diagonal 2×2 matrices with entries v_1, v_2 and φ_1, φ_2, respectively. $g_{L,R}$ represent Keldysh matrices of the electron propagators in the leads at the position of the dot. Their Fourier transforms are known: $g^r = -i\pi$, $g^{12}(\epsilon) = 2i\pi f(\epsilon)$ with $f(\epsilon)$ as Fermi distribution function. All dependencies on constant bias and fluctuating potentials are collected in the exponents of Eq.(10). We also define $T_k = \Gamma b^2$ as an effective Kondo temperature. The $G(\omega)$ includes the Lagrange multiplier which shifts the localized level position ϵ to $\tilde{\epsilon} = \epsilon + \lambda$. For our purpose it is enough to find both free parameters, T_k and renormalized level $\tilde{\epsilon}$, by solving two self-consistent equations in equilibrium (neglecting the fluctuations of electrical potential) [9]. These equations follow from extremum conditions of effective action S_{eff} relative to b and $\tilde{\epsilon}$

$$\frac{\partial S_{eff}}{\partial b} = \frac{\partial S_{eff}}{\partial \tilde{\epsilon}} = 0 \tag{13}$$

Unlike the multichannel Kondo model, the solution for spin 1/2 Kondo quantum dot satisfies [9] $\tilde{\epsilon} \sim T_K(T_K/\Gamma) << T_K$ which permits us to use $\tilde{\epsilon}/T_K$ as a small parameter.

To obtain the effective action at one loop level we have to expand the logarithm in equation (8) to the second power in quantum field χ. The calculations are straightforward, though, rather long. The results for the unitary limit are in accordance with those obtained by bosonization technique [6, 11, 12, 13] and in reference [4]. The part of effective action ($\sim Ln$) (8) which describes the electron scattering thus acquires a form

$$\begin{aligned} \tilde{S} &= \frac{N}{\pi}[\frac{i}{2}\int\int dtdt'\chi(t)\gamma(t-t')\chi(t') \\ &- \int dt\chi(t)(\dot{\varphi}_1+\dot{\varphi}_2)/2] \end{aligned} \tag{14}$$

Here the number of the channels $N = 2$. Collecting together all terms we arrive to a compact form of effective action which is the principal result of this section:

$$S_{eff} = \int\int dtdt'\hat{\varphi}(t)\mathbf{B}^{-1}(t-t')\hat{\varphi}(t') \tag{15}$$

The Fourier image of Coulomb Green function $\mathbf{B}(t)$ we find $\mathbf{B}(\omega) = \mathbf{B}_v(\omega)/\omega^2$ where

$$\mathbf{B}_v(\omega) = \frac{E_c}{\omega^2+Q^2}\begin{pmatrix} \omega^2 - iQ\gamma & iQ(\omega-\gamma) \\ -iQ(\omega+\gamma) & -\omega^2 - iQ\gamma \end{pmatrix} \tag{16}$$

here $Q = E_c\bar{N}/\pi$, $\bar{N} = N(1+h/(4e^2R_s))$ is the effective number of conducting channels, and $\gamma \equiv \gamma(\omega)$. $\mathbf{B}_v(\omega)$ is matrix Green function (GF) of electric potentials v_i. Having the effective action we can calculate correlation functions of fluctuating potential fields and, thus, find the conductance and shot noise corrections in the Fermi liquid regime of the Kondo model.

Fermi Liquid Regime: Near the unitary limit it is convenient to follow the scattering approach. For this purpose we use the basis of s and p scattering states rather than those of the left-lead and right-lead states. At $V = 0$ the p states are decoupled from the dot. For nonzero fluctuating potentials the p-states are not decoupled, however, the s, p basis is important. In real junctions with strong Kondo effect, nearly $T = 0$ the dot is described by strongly interacting fix-point with many body state ($\bar{b}$). Before scattering the phases of both states, $\bar{b}$ and s -state, coincide, while passing the scattering region $\bar{b}$ state acquires only an extra phase π compared to that of s-state. Then the fix-point Hamiltonian can be written [8, 14, 15] in the new basis $(b_{\sigma,k}, a_{\sigma,k})$ (here $a_{\sigma,k}$ stands for p-states) as $H = H_0 + H_s + H_{int}$

$$\begin{aligned} H_0 &= \sum_{k,\sigma}\xi_k(b^\dagger_{\sigma,k}b_{\sigma,k} + a^\dagger_{\sigma,k}a_{\sigma,k}) \\ &+ \frac{v}{2}\sum_{k,\sigma}(b^\dagger_{\sigma,k}a_{\sigma,k} + a^\dagger_{\sigma,k}b_{\sigma,k}) \end{aligned}$$

$$H_s = -\frac{a}{\nu T_K}\sum_{k,k',\sigma}(\xi_k+\xi_{k'})b^\dagger_{\sigma,k}b_{\sigma,k'} \tag{17}$$

$$H_{int} = \frac{b}{\nu^2 T_K}b^\dagger_\uparrow b_\uparrow b^\dagger_\downarrow b_\downarrow \tag{18}$$

where $b = 2a$ and $a = 1/(2\pi)$ [14] and $b_\uparrow = \sum_k b_{\uparrow,k}$. The Kondo temperature T_K is the only energy scale of the fixed-point Hamiltonian. We can integrate out the $a_{\sigma,k}$-states and obtain the effective action on Keldysh contour only for interacting states $\bar{b}$

$$\begin{aligned} S &= \int\int dtdt'\sum_{k\sigma}b^\dagger_{k\sigma}(t)G^{-1}_{bk}(t-t')b_{k\sigma}(t') \\ &- \int dt(H^i_s + H^i_{int})\sigma^{ii}_z \end{aligned} \tag{19}$$

$$G^{-1}_{bk}(t,t') = [(i\frac{\partial}{\partial t} - \xi_k)\sigma_z - \frac{1}{4}\hat{v}\sigma_z G_{ak}\sigma_z\hat{v}] \tag{20}$$

Here G_{ak} is matrix Green function for decoupled noninteracting a- states (for example, $G^r_{ak}(\epsilon) = 1/(\epsilon - \xi_k + i\delta)$). We separate potentials in (20) into fluctuating part and constant applied voltage V, that is, write $\hat{v} \to V + \hat{v}$. To obtain the Coulomb interaction corrections to conductance and shot noise we have to express the current in terms of Green functions (the inverse of (20) and then, with help of effective action (15), to average the product of fluctuating potentials. This goal can be achieved by applying the perturbation theory (it corresponds to 1/N expansion).

Current: To begin with, we consider, at first, the average current. The backscattering current operator is given as [1, 2, 3, 8]:

$$\begin{aligned} I_b &= -\frac{ie}{\hbar}\frac{a}{2\nu T_K}\sum_{k,k',\sigma}(\xi_k+\xi_{k'})(b^\dagger_{\sigma,k'}a_{\sigma,k} - a^\dagger_{\sigma,k}b_{\sigma,k'}) \\ &- \frac{ie}{\hbar}\frac{b}{2\nu^2 T_K}\sum_\sigma(b^\dagger_\sigma a_\sigma - a^\dagger_\sigma b_\sigma)n_{\bar{\sigma}} \end{aligned} \tag{21}$$

and $n_\sigma = b^\dagger_\sigma b_\sigma$, $\bar{\sigma} = -\sigma$. The sum over σ stands for spin summation, which in our case is trivial. The averaged backscattering current consists of two parts $I_b = I_s + I_{int}$, where I_s and I_{int} originate, respectively, from the scattering Hamiltonian H_s and the interaction H_{int}. I_s and I_{int} are generally expressed in terms of Green's functions for a and $\bar{b}$ states.

$$I_s = -\frac{e}{\hbar}(\frac{a}{2\nu T_K})^2\sum_{k,k',\sigma}(\xi_k+\xi_{k'})^2(A_{k,k'} - \bar{A}_{k,k'}) \tag{22}$$

$$I_{int} = -\frac{e}{\hbar}(\frac{b}{2\nu T_K})^2\sum_{k,\sigma}(\Pi_k - \bar{\Pi}_k) \tag{23}$$

here

$$\begin{aligned} A_{k,k_1} &= <[G_{ak}\hat{v}\sigma_z G_{bk}\sigma_z G_{bk_1}]^{11}_{tt}> \\ \bar{A}_{k,k_1} &= <[G_{bk_1}\sigma_z G_{bk}\hat{v}\sigma_z G_{ak}]^{11}_{tt}> \end{aligned} \tag{24}$$

$$\begin{aligned} \Pi_k &= <[G_{ak}\hat{v}\sigma_z G_{bk}\sigma_z G_b]^{11}_{tt}[G_b G_b]^{11}_{tt}> \\ \bar{\Pi}_k &= <[G_b\sigma_z G_{bk}\hat{v}\sigma_z G_{ak}]^{11}_{tt}[G_b G_b]^{11}_{tt}> \end{aligned} \tag{25}$$

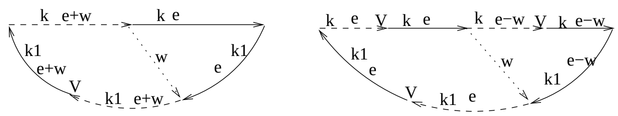

FIG. 2: Diagrams for processes which contribute to I_s. The solid lines with momentum k denotes $\bar{b}$-electron Green functions $g_{bk}(\epsilon)$. Dashed lines represent GF of noninteracting states a. The dot line stands for Coulomb GF $iB_v(w)$ (16). All Green functions are 2×2 Keldysh GF. Bias voltage V denotes the vertex that connects dash and solid line. Each vertex (but most left one) includes σ_z, and matrix product is considered.

where brackets $<$ ... $>$ denote the average over fluctuating potentials with the effective action (15). The expression in square brackets is matrix product in Keldysh space and convolution in time. One has to take the 11 component of this matrix product. The last expression (23), unlike the scattering part of the current (22), includes, in addition, the Green functions at the dot position $G_b = \sum_k G_{bk}$. In the zero order approximation (no potential fluctuations) for the Fermi liquid regime of Kondo effect the nonlinear current was found before [8, 14, 15] (see also recent works [1, 2, 3], where the explicit separation between scattering and interacting contributions to conductance was done). In this case the backscattering current reads as,

$$I_b^0 = \frac{2e^2V}{3h}[\pi^2(a^2+b^2\frac{5}{4})(\frac{eV}{T_K})^2] \qquad (26)$$

Let us now include the fluctuations of potential fields. At first we do the calculations for I_s. For this purpose we expand G_b to the second order in fluctuating electrical potentials. The GFs for $\bar{b}$-states to zero approximation in fluctuating fields acquire simple form:

$$\begin{aligned}
g^R_{bk}(\omega) &= \frac{\pi i}{2}(\frac{1}{\omega-\xi_k-V/2+i\delta}+\frac{1}{\omega-\xi_k+V/2+i\delta}] \\
g^{12}_{bk}(\omega) &= -f(\xi_k)(g^R(k\omega)-g^A(k\omega)) \\
g^{21}_{bk}(\omega) &= (1-f(\xi_k))(g^R(k\omega)-g^A(k\omega)) \\
g^{i,j}_{b}(\omega) &= \int d\xi_k g^{i,j}_{bk}(\omega)
\end{aligned}$$

The diagrams are a convenient way to represent different contributions to the transport current. For I_s two types of diagram are shown in Fig2. The principal contributions to the scattering part of the current at $1/N$ order are given by a number of processes. These processes are presented by: (i) two diagrams in Fig2, (ii) two diagrams similar to those in Fig2, however, with the Coulomb (dot) line connected with other vertex (V) (these diagrams are obtained by an interchange between the V vertex and the end of the dot line), and, (iii) four "conjugate" diagrams for $\bar{A}_{k,k'}$ (like the cases (i) and (ii)). These "conjugate" diagrams together with diagrams in Fig2 simply double the real parts of those in Fig2. We notice that there are yet processes which can be find by self-energy insertion in the one momentum line. However, their impact on current is small. Thus summing all relevant contributions and including the part without fluctuations one finds the scattering part of the total current

$$I_s = \frac{2e^2Va^2\pi^2}{3h}(\frac{eV}{T_K})^2[1+\frac{9}{16\bar{N}}(\ln[\frac{2Q}{eV}]+1)] \qquad (27)$$

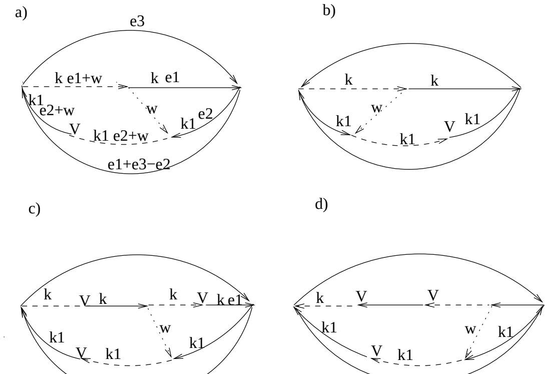

FIG. 3: The processes that contribute to the interacting part of transport current I_{int}. The notations for momentum dependent GF are the same as in Fig1. The lines without momentum denote integrated over momentum GFs g_b.

We need more diagrams (see Fig3) to calculate I_{int}. Diagrams of types (a)and (c) are natural generalization of those in Fig2. Diagrams of type (b) and (d) are of "scattering" type and appear only due to the interactions. The accounting for relevant processes goes as follow: at first the diagrams for $\bar{\Pi}$ are "conjugate to those of Π and after summation, as in the case for I_s, they double the real parts of diagrams in Fig3. Next, we have :(i) four diagram of type (a)(two of them appear due to interchange with vertex V), (ii) $2\bigotimes 4 = 8$ diagram similar to (c)(factor four appears due to four ways to connect V vertices), (iii) two diagrams of type (b), and (iv) four diagrams of type (d). The summation all of this contributions results in the Coulomb interaction correction to the I_{int}

$$I_{int} = \frac{2e^2Vb^2\pi^2}{3h}\frac{5}{4}(\frac{eV}{T_K})^2[1+\frac{2}{\bar{N}}(\ln[\frac{2Q}{eV}]+0.11)] \qquad (28)$$

The corresponding corrections to differential conductance due to electron-electron interaction are obtained by taking the derivative of (27), (28) with respect to V.

Shot Noise: The effective charge of the current-carrying particles is defined as $e^* = S_n/I_b$, where S_n is the zero frequency Fourier transform of current-current

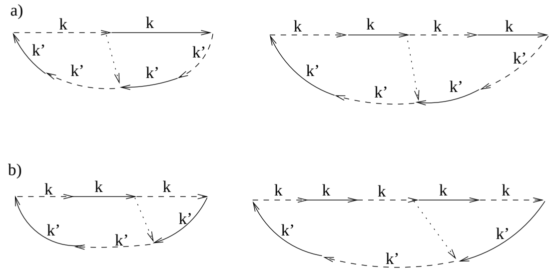

FIG. 4: Diagrams which contribute into the scattering part of the noise. Notations of the GFs and momentums are the same as in Fig2, however, the time contour indices of the edge vertices are different: here index 2 is in the left and index 1 in the right vertex. Diagrams b) enter with sign minus.

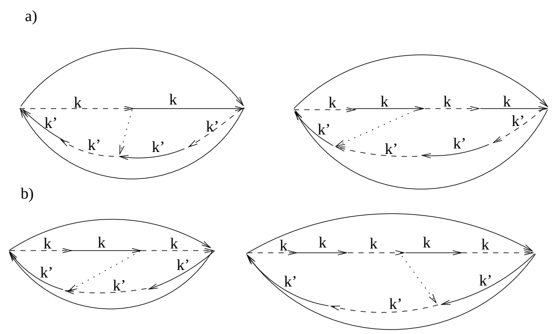

FIG. 5: The processes that contribute to the interacting part of shot noise S_n^{int}. The notations for GF correspond to those in Fig2 and Fig3. The edge Keldysh vertex indices are the same as in Fig3. Diagrams b) enter with minus sign.

correlation function,

$$S(t) \equiv < I(t)I(0) > - < I >^2, \tag{29}$$

The same as the current, noise power can be separated into the two terms $S = S_n^s + S_n^{int}$ related, respectively, to the scattering part of the Hamiltonian H_s and to the interaction part H_{int}. In the Fermi liquid regime (without fluctuating potentials) at temperature $T = 0$ the shot noise was obtained before [1, 3]

$$S_n^0 = \frac{e^2 V}{h}\pi^2(\frac{3}{2}b^2 + \frac{2a^2}{3})\frac{(eV)^2}{T_K^2} \tag{30}$$

The influence of Coulomb interaction changes the scattering and interacting parts of S_n. Explicitly, we can find these changes (like in the case of the current) at $1/N$ order of perturbation theory. In the formula for noise (30) we assign Keldysh indices to the current operators (21), and then, using Keldysh technique, sum the all relevant diagrams. For scattering part S_n^s these diagrams are shown in Fig.4. There are also equivalent (not represented by Fig.4) diagrams. They also contribute to the noise, and their number, like for current, can be easily counted. Thus, for the total scattering noise power which includes potential fluctuations, we obtain

$$S_n^s = \frac{2e^2 V a^2 \pi^2}{3h}(\frac{eV}{T_K})^2[1 + \frac{1}{2\bar{N}}(\ln[\frac{2Q}{eV}] + 1.18)] \tag{31}$$

For interacting part of the shot noise S_n^{int} the typical diagrams are represented in Fig.5. In addition to them there are also "scattering" diagrams (GFs of inside bobble point in one direction). Each such diagram corresponds (with the same contribution) to one in Fig.5. After summation of all diagrams we obtain

$$S_n^{int} = \frac{3e^2 V b^2 \pi^2}{2h}(\frac{eV}{T_K})^2[1 + \frac{5}{6\bar{N}}(\ln[\frac{2Q}{eV}] + 0.31)] \tag{32}$$

By using Eqs.(31,32,27,28) and definition of the e^* we can immediately obtain factor F (see Eq.1).

Conclusions: We have presented a theory to evaluate the effects of electron-electron interactions on quantum transport through the Kondo quantum dot. We address the important problem of effective backscattering charge for current-carrying particles in the interacting systems. This charge can be obtained from the measurements of shot noise and backscattering current in the Fermi liquid regime of a Kondo quantum dot. We show that in this regime the potential fluctuations can markable reduce the recently calculated universal value $5/3e$ [1], though, the effective charge remains bigger then e. To get our principal result we have calculated the transport current and obtained the zero frequency shot noise power. Of cause, the potential fluctuations are not the only cause to change the effective charge. The pure unitary limit in real system may by broken by the symmetry breaking interactions like potential scattering or magnetic field which can also influence the effective charge. However, in the case these effects are small, they can be properly estimated [1, 3, 14, 15] and render a small correction to e^*.

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